

Large loop effects of extra SUSY Higgs doublets to CP violation in B^0 mixing

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Abstract

We consider more than one pair of $SU(2)_L$ doublet Higgs supermultiplets in a generic supersymmetric extension of the standard model, and calculate their one-loop contributions to the soft mass insertions δ_{LL} etc. We find that if large supersymmetry breaking in this sector is realized, the loop effects can give rise to large contributions to the soft mass insertions, meaning that they can generate large FCNCs and CP violations.

We apply our result to a recently proposed model based on the discrete Q_6 family group, and calculate the non-diagonal matrix element M_{12} of the neutral meson systems. We focus our attention on the extra phases $\phi_{d,s}^\Delta$ in $B_{d,s}$ -mixing and flavor-specific CP-asymmetries $a_{sl}^{d,s}$ in neutral B decays and obtain values that can be about one order of magnitude larger than the standard model predictions. Our final results are comparable with the recent experimental observations at D0 and CDF, but they are still about a factor of 5 smaller than the recently measured dimuon asymmetry from D0.

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I. INTRODUCTION

The CKM mechanism [1] has been tested and confirmed to a high accuracy as the dominant source of flavor violation and CP violation in the standard model (SM), see e.g. [2, 3]. Despite this success it is well-known that the amount of CP violation present in the SM is not sufficient [4] to explain the baryon asymmetry in the universe [5]. Moreover in the last years also some hints for deviations of the CKM picture were accumulated. In particular the recent measurement of the CP-violating dimuon asymmetry by the D0 collaboration [6] gained a lot of attention. The measured value

$$A_{sl}^b = -(9.57 \pm 2.51 \pm 1.46) \cdot 10^{-3}, \quad (1)$$

is a factor of 42 larger than the SM prediction [7]

$$A_{sl}^b = -(2.3_{-0.6}^{+0.5}) \cdot 10^{-4}. \quad (2)$$

The statistical significance of this deviation is 3.2σ . Since this large deviation might hint to a sizeable new source of CP violation, needed to solve the problem of the baryon asymmetry, the D0 measurement resulted already in many theoretical papers investigating different new physics models [8], for earlier works on large CP violation in B_s mixing see e.g. [9–11].

The SM and its minimal supersymmetric extension (MSSM) with softly broken supersymmetry have the minimal structure of the Higgs sector. In the SM this minimality is the main reason that the flavor-changing-neutral current (FCNC) as well as the CP-violating processes are highly suppressed. However, the minimality of the Higgs sector in the MSSM does not help suppressing FCNC and CP violation at all, and this gives arise to the well-known SUSY flavor and CP problem [12–20]. Therefore, one is led to consider another mechanism to suppress FCNC and CP violation in supersymmetric extensions of the SM. A natural assumption is that spontaneous CP violation or its modification is responsible for the small CP violation in the MSSM. However, spontaneous CP violation in the MSSM does not occur, unless one extends the Higgs sector to a non-minimal form [21]. Moreover, just adding more $SU(2)_L$ doublet Higgs supermultiplets does not help; one should introduce a certain set of SM singlet Higgs bosons [21]. However, an extension to introduce more than one pair of $SU(2)$ doublet Higgs supermultiplets might suffer from two major problems:

- (i) it can destroy the successful gauge coupling unification, and
- (ii) large tree-level FCNCs can be present.

We will ignore problem (i) in this work, while we will solve problem (ii) by introducing a flavor symmetry; a symmetry-based mechanism to suppress FCNCs was considered in the literature e.g. in [22]–[27], [10, 11]¹. Note, however, that if CP violation in B^0 mixing should

¹ Recently, a considerable attention has been given to the idea of incorporating a non-abelian flavor symmetry into a GUT[28].

turn out to be large as the D0 measurement (1) is suggesting, we are caught in a dilemma between suppressed and large CP violation. In this paper we would like to address this dilemma (see also S. King of [8], who has also addressed the problem in a similar framework very recently).

First we consider more than one pair of $SU(2)_L$ doublet Higgs supermultiplets in a generic supersymmetric extension of the SM. We then calculate the one-loop contributions of the extra heavy Higgs multiplets to the soft mass insertions δ 's [12, 13]. We find that the loop effects can give rise to large contributions to the soft mass insertions. That is, the loop effects can generate large FCNCs and CP violations in such models. We then apply our general result to a specific model, in which the problem (ii) is overcome by the flavor symmetry Q_6 [25], and investigate the one-loop effects mentioned above on the dimuon asymmetry and related observables. The Higgs sector of this model consists of six $SU(2)_L$ doublets, three for the up-quark sector and three for the down-quark sector; the three doublets of each sector form a three-dimensional reducible representation $\mathbf{1} + \mathbf{2}$ of Q_6 . In the end four of the six $SU(2)_L$ doublets are super-heavy \gtrsim few TeV (which comes from the FCNC constraints in the mixing of the neutral meson systems [11, 27]), and two of them form the pair of the MSSM Higgs supermultiplets. Tree-level contributions to the semileptonic asymmetries, due to the exchange of the extra heavy neutral Higgs bosons, were discussed in [10]. There it was found that the small standard model expectations for semileptonic CP asymmetries can be enhanced by up to one order of magnitude. In this paper we assume that the extra heavy Higgs bosons are so heavy that the tree-level contributions can be neglected, and that the extra FCNCs and CP violations come only from the SUSY breaking sector. We determine under this assumption in this model the non-diagonal matrix element M_{12} of the neutral meson systems and investigate the possible size of a new CP-violating phase as well as the possible size of semileptonic CP-asymmetries. Our results are finally compared with the recent measurements, in particular the dimuon asymmetry of D0 [6].

II. NEW PHYSICS IN B^0 MIXING

The mixing of neutral mesons is governed by the famous box diagrams. The dispersive part, denoted by M_{12} is expected to be very sensitive to new physics, while for the absorptive part, denoted by Γ_{12} , new contributions are expected to be below the hadronic uncertainties². Therefore one can write generally [7] in the presence of new physics

$$M_{12}^q = M_{12}^{SM,q} \cdot \Delta_q, \quad \Delta_q = |\Delta_q| e^{i\phi_q^\Delta}, \quad (3)$$

$$\Gamma_{12}^q = \Gamma_{12}^{SM,q}. \quad (4)$$

² This statement was recently questioned in the literature. A more detailed discussion of it, will be given in [29].

$q = s$ denotes the B_s -system and $q = d$ denotes the B_d -system. Defining the phase ϕ_q as

$$\phi_q = \arg\left(-\frac{M_{12}^q}{\Gamma_{12}^q}\right), \quad (5)$$

we get [7] the following general expression for the flavor-specific CP-asymmetries (sometimes also called semileptonic CP-asymmetries) in the presence of new physics

$$a_{sl}^q = \text{Im}\left(\frac{\Gamma_{12}^q}{M_{12}^q}\right) = \frac{|\Gamma_{12}^q|}{|M_{12}^{SM,q}|} \cdot \frac{\sin(\phi_q^{SM} + \phi_q^\Delta)}{|\Delta_q|}. \quad (6)$$

The above formula holds independent of the assumption, that there is almost no new physics possible in Γ_{12} . In the SM the CP-violating phases ϕ_q and the semileptonic CP-asymmetries are small, one gets [7] (based on the results of [30–34])

$$a_{sl}^d = (-4.8_{-1.2}^{+1.0}) \cdot 10^{-4}, \quad \phi_d^{SM} = -0.091_{-0.038}^{+0.026} = -5.2_{-2.1}^{+1.5}^\circ, \quad (7)$$

$$a_{sl}^s = (2.06 \pm 0.57) \cdot 10^{-5}, \quad \phi_s^{SM} = (4.2 \pm 1.4) \cdot 10^{-3} = 0.24^\circ \pm 0.08^\circ. \quad (8)$$

D0 measured [6] a linear combination of the semileptonic CP-asymmetries in the B_d and in the B_s system

$$A_{sl}^b = (0.494 \pm 0.043) \cdot a_{sl}^s + (0.506 \pm 0.043) \cdot a_{sl}^d. \quad (9)$$

The experimental central value turned out to be a factor of 42 larger than the SM expectation for A_{sl}^b . Using the experimental value for $a_{sl}^d = -0.0047 \pm 0.0046$ from [35] one derives [6] a bound on a_{sl}^s :

$$a_{sl}^s = (-14.6 \pm 7.5) \cdot 10^{-3}. \quad (10)$$

Inserting this value in Eq.(6) we get with the results from [7]

$$\sin(\phi_s^{SM} + \phi_s^\Delta) = -(2.9 \pm 1.5) \cdot |\Delta_s|. \quad (11)$$

Assuming however, that there is no new physics in B_d -mixing one gets instead

$$a_{sl}^s = (-19 \pm 10) \cdot 10^{-3} \Rightarrow \sin(\phi_s^{SM} + \phi_s^\Delta) = -(3.8 \pm 2.0) \cdot |\Delta_s|. \quad (12)$$

Using the fact that $|\Delta|$ is closed to one (see e.g. [7]) we get in both cases unphysical values for $\sin(\phi_s^{SM} + \phi_s^\Delta)$. This problem will be discussed in detail in [29], here we simply assume that the current data hint for a large value of ϕ_s^Δ compared to the SM angle ϕ_s^{SM} . This also holds if we combine the D0 dimuon asymmetry with previous direct determinations of semileptonic CP-asymmetries [36].

If a non-vanishing value of ϕ_s^Δ is realized in nature, this would also be visible in the angular analysis of the decay $B_s \rightarrow J/\Psi\Phi$ [37]. In the SM one extracts in this decay the angle $-2\beta_s \approx -2.2^\circ$ (for the notation see e.g. Noted added in [38]). If new physics is only present in the B_s -mixing and not in the $B_s \rightarrow J/\Psi\Phi$ decay one extracts instead the angle $-2\beta_s + \phi_s^\Delta$.

Current data [39–41] for $B_s \rightarrow J/\Psi\Phi$ also hint to a non-vanishing value of $\phi_s^{\Delta 3}$, which points to the same direction as the value of the semileptonic CP-asymmetries measured by D0 [6]. Possible problems using this extraction for the CP-violating phase in B_s mixing are discussed in detail in [29].

III. GENERAL FORMULA

Consider the superpotential

$$W = \hat{Y}_{ij}^{uI} Q_i U_j^c H_I^u + \hat{Y}_{ij}^{dI} Q_i D_j^c H_I^d + \mu^{IJ} H_I^u H_J^d. \quad (13)$$

Here $SU(2)_L$ doublets of the quark and Higgs supermultiplets are denoted by Q, H^u and H^d , respectively. The indices I and J indicate different kinds of the Higgs $SU(2)_L$ doublets. Similarly, U^c and D^c stand for $SU(2)_L$ singlets of the quark supermultiplets. We denote the component fields by a small letter along with a \sim for the scalar quarks and higgsinos, and $(\hat{Y})^* = Y$. (Y is the Yukawa coupling in our notation, i.e. $Y \bar{q}_L q_R$.) The $SU(2)$ components of the Higgs fields are

$$h_I^u = (\phi_I^{u+}, \phi_I^{u0}), \quad h_I^d = (\phi_I^{d0}, \phi_I^{d-}). \quad (14)$$

To compute the corrections to the soft mass insertions $(\delta_{ij})_{LL, \text{etc}}$, which will be defined in (50), we have to compute the corrections to the squark masses. In the following calculations we consider only the insertions $(\delta_{12}^d)_{LL, LR}$ for the down-type scalar quarks, and moreover neglect the Y^d 's except for the tree-level δ_{LR} . Quark and squark masses are also neglected. Calculation of other types of the insertions can be done in a similar way. There will be (i) tree-level contributions to δ_{LR} coming from the fact that there are more than one pair of Higgs doublets. Then (ii) diagrams of Fig. 1 with heavy Higgs fields in the loop, and two types of loop diagrams; (iii) those with the heavy Higgs bosons and squarks in the loop as shown in Fig. 2 (a) and (vi) those with the higgsinos and quarks in the loop as shown in Fig. 2 (b). We obtain the following results:

(i) Tree-level contribution to $(\delta_{ij}^d)_{LR}$:

The relevant Lagrangian is

$$\mathcal{L}_\mu = - \left(\phi_I^{u0} \mu^{IJ} Y_{ij}^{dJ} \right) \tilde{d}_{Li}^* \tilde{d}_{Rj} + h.c., \quad (15)$$

which yields

$$(\delta_{ij}^d)_{LR}(\mu) = \left(\frac{\langle \phi_I^{u0} \rangle \mu^{IJ}}{m_d^2} \right) \left[(U_L^d)^\dagger Y^{dJ} U_R^d \right]_{ij}, \quad (16)$$

³ A new result from CDF [42] was presented at FPCP 2010 giving a 1-sigma range of $\phi_s^\Delta \in [0, -1]$, being perfectly consistent with the SM ($\phi_s^\Delta = 0$), but also with a large deviation from the SM ($\phi_s^\Delta = -1 \approx -57^\circ$).

where $m_{\tilde{d}}$ is the average squark mass, and U 's are unitary matrices that diagonalize the down-type quark mass matrix, i.e.

$$(U_L^d)^\dagger \mathbf{m}_d U_R^d = \text{diag.}(m_d, m_s, m_b) . \quad (17)$$

There are besides the usual contribution to δ_{LR} coming from the A terms, which we have not included. The contribution (16) exists only if there are more than one pair of Higgs doublets.

(ii) Quartic coupling contribution to $(\delta_{ij}^d)_{LL}$ (Fig. 1):

The quartic couplings are given by

$$\mathcal{L}_{\text{quart}} = -Y_{ik}^{uI} (Y^{uJ})_{kj}^\dagger (\phi_I^{u+})^* \phi_J^{u+} \tilde{d}_{Li}^* \tilde{d}_{Lj} . \quad (18)$$

We find

$$(\delta_{ij}^d)_{LL}(\text{quart}) = \frac{[(U_L^d)^\dagger Y^{uI} (Y^{uJ})^\dagger U_L^d]_{ij}}{16\pi^2 m_d^2} (U_c \mathbf{M}_c^{2,dia} U_c^\dagger)_{JI} , \quad (19)$$

where

$$\phi_I^{u+} = U_{c,IJ} \phi_J^{u+,dia} , (\mathbf{M}_c^{2,dia})_{IJ} = (m_{c,J}^2 \ln m_{c,J}^2 / Q^2) \delta_{IJ} , U_c^\dagger \mathbf{M}_c^2 U_c = m_{c,J}^2 \delta_{IJ} . \quad (20)$$

\mathbf{M}_c^2 is the mass matrix for the charged Higgs bosons, and Q is the renormalization scale. We have suppressed the constant terms which, however, can be absorbed into Q . We will do so for other diagrams.

(iii) Cubic scalar couplings to $(\delta_{ij}^d)_{LL}$ (Fig. 2 (a)):

$$\mathcal{L}_{\text{CS}} = Y_{ij}^{uI} \mu^{IJ} \phi_J^{d-} \tilde{d}_{Li}^* \tilde{u}_{Rj} + h.c. , \quad (21)$$

which gives

$$(\delta_{ij}^d)_{LL}(\text{CS}) = \frac{[(U_L^d)^\dagger Y^{uI} (Y^{uJ})^\dagger U_L^d]_{ij}}{16\pi^2 m_d^2} (\mu U_c^* \mathbf{L}_c^{2,dia} U_c^T \mu^\dagger)_{JI} , \quad (22)$$

where $(\mathbf{L}_c^{2,dia})_{IJ} = (\ln m_{c,J}^2 / Q^2) \delta_{IJ}$, and U_c and $m_{c,J}$ are defined in (20). Here we have neglected the mass of \tilde{u}_R .

(vi) Higgsino loop to $(\delta_{ij}^d)_{LL}$ (Fig. 2 (b)):

$$\mathcal{L}_h = -Y_{ij}^{uI} \tilde{d}_{Li}^* \tilde{h}_I^{u+} u_{Rj} + h.c. \quad (23)$$

The expression for $(\delta_{ij}^d)_{LL}(h)$ is similar to the quartic coupling contribution (19). We find, neglecting the quark masses,

$$(\delta_{ij}^d)_{LL}(h) = -2 \frac{[(U_L^d)^\dagger Y^{uI} (Y^{uJ})^\dagger U_L^d]_{ij}}{16\pi^2 m_d^2} [U_h \mathbf{M}_h^{2F,dia} U_h^\dagger]_{JI} , \quad (24)$$

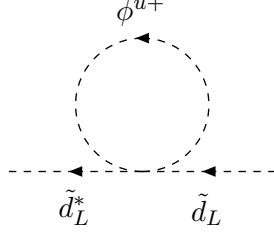


FIG. 1: The quartic coupling contribution to $(\delta_{ij}^d)_{LL}$.

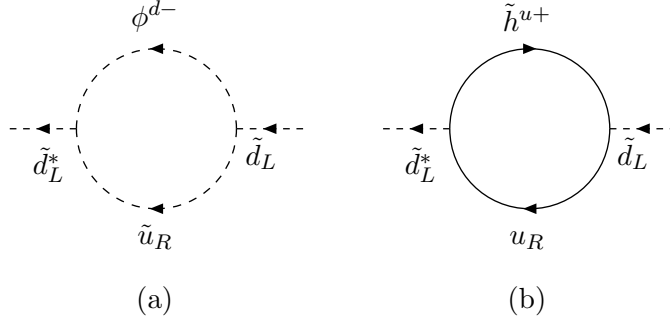


FIG. 2: The heavy Higgs (a) and higgsino (b) contributions to $(\delta_{ij}^d)_{LL}$.

where $\tilde{h}_I^{u+} = U_{h,IJ} \tilde{h}_J^{u+,dia}$, and

$$(\mathbf{M}_h^{2F,dia})_{IJ} = (m_{h,J}^2 \ln m_{h,J}^2/Q^2) \delta_{IJ} , \quad (U_h^\dagger \mathbf{M}_h^F U_h)_{IJ} = m_{h,J} \delta_{IJ} . \quad (25)$$

\mathbf{M}_h^F is the mass matrix for the charged higgsinos.

Before we apply the results above we make few remarks. The infinite renormalization of the soft scalar masses do not depend on the μ and A terms to all orders in perturbation theory [43]. Therefore, the μ parameter dependence of the infinite part (and hence of $\ln Q$) in δ 's should be cancelled. However, the cancellation of the finite part is not exact. As we see from (22) - (24), the insertions $(\delta)_{LL}$'s explicitly depend on μ parameters (\mathbf{M}_c^2 and \mathbf{M}_h^F also contain μ parameters). Keeping this in mind, we consider

$$D = \mu_1^2 \ln m_1^2/Q^2 + \mu_2^2 \ln m_2^2/Q^2 - (\mu_1^2 + \mu_2^2) \ln m_3^2/Q^2 \quad (26)$$

in which the renormalization scale Q dependence exactly cancels. If all m_i 's are of the same size, D is small compared to the μ^2 's. However, if there is a large SUSY breaking so that the mass of a fermionic component (higgsino) differs from that of the bosonic component (Higgs) by a large amount, D may become large. Moreover, there are terms in δ 's which, instead of μ^2 , are proportional to the square of the soft scalar masses of the Higgs bosons,

	Q	Q_3	U^c, D^c	U_3^c, D_3^c	L	L_3	E^c, N^c	E_3^c	N_3^c	H^u, H^d	H_3^u, H_3^d
Q_6	$\mathbf{2}_1$	$\mathbf{1}_{+,2}$	$\mathbf{2}_2$	$\mathbf{1}_{-,1}$	$\mathbf{2}_2$	$\mathbf{1}_{+,0}$	$\mathbf{2}_2$	$\mathbf{1}_{+,0}$	$\mathbf{1}_{-,3}$	$\mathbf{2}_2$	$\mathbf{1}_{-,1}$

TABLE I: The Q_6 assignment of the chiral matter supermultiplets, where the group theory notation is given in Ref. [25]. For completeness we include leptons, L, E^c and N^c . R parity is also imposed.

which we denote generically by m_s . If $m_s^2 \gg \mu^2$, these m_s^2 terms dominate. Then, there will be corrections $\sim y^2(m_s^2/16\pi^2 m_d^2) \ln m_s^2/Q^2$ to the δ 's, where y stands for a generic Yukawa coupling. So, if $m_s/m_d \gg 1$ is realized by one reason or another, the loop effects of the heavy Higgs bosons to the soft mass insertions δ 's may become large. One reason may be the following. Flavor-changing neutral Higgs bosons should be made heavy to suppress tree-level FCNC processes, while higgsinos should be light because we need small μ 's to suppress EDMs which are caused by the tree-level $(\delta)_{LR}$ given in (16), as was pointed out in [11]. To make the Higgs bosons heavy, we have to make the corresponding soft scalar masses large, but not the μ parameters from the reason above. This leads to a large SUSY breaking.

IV. AN APPLICATION

In this section we would like to apply our results of the last section to a specific model. The model [25] is based on a discrete flavor symmetry along with spontaneous CP violation. In this model the tree-level CP violation is suppressed [10, 11], and it is difficult to explain possible large CP violations, for which recently some evidence was observed at D0 [6]. We would like to see whether the heavy Higgs fields present in the model can help obtaining large CP violation in the $B^0 - \bar{B}^0$ mixing.

A. The model

The model is briefly described below (the details of the model can be found in [11, 26, 27]). The Q_6 assignment is shown in Table I. Here we restrict ourselves to the quark sector. The most general Q_6 invariant, renormalizable superpotential for the Yukawa interactions in the quark sector yields the following Yukawa matrices [25]:

$$\begin{aligned}
\mathbf{Y}^{u1(d1)} &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & Y_b^{u(d)} \\ 0 & Y_{b'}^{u(d)} & 0 \end{pmatrix}, \quad \mathbf{Y}^{u2(d2)} = \begin{pmatrix} 0 & 0 & Y_b^{u(d)} \\ 0 & 0 & 0 \\ -Y_{b'}^{u(d)} & 0 & 0 \end{pmatrix}, \\
\mathbf{Y}^{u3(d3)} &= \begin{pmatrix} 0 & Y_c^{u(d)} & 0 \\ Y_c^{u(d)} & 0 & 0 \\ 0 & 0 & Y_a^{u(d)} \end{pmatrix}.
\end{aligned} \tag{27}$$

All the Yukawa couplings are real, because we assume that CP is spontaneously broken. So, the VEVs of the Higgs fields have to be complex to obtain the CP phase of the CKM matrix. It has been found in [25] that for successful predictions, i.e. sum rules among the fermion masses and CKM parameters, the scalar potential should have an accidental Z_2 invariance

$$h_+^{u,d} = \frac{1}{\sqrt{2}}(h_1^{u,d} + h_2^{u,d}) \rightarrow h_+^{u,d}, \quad h_-^{u,d} = \frac{1}{\sqrt{2}}(h_1^{u,d} - h_2^{u,d}) \rightarrow -h_-^{u,d}. \quad (28)$$

Then there will be only nine independent parameters in the quark sector to describe ten observables (six quark masses and four CKM parameters).

The superpotential for the Higgs sector consists of μ terms. The only Q_6 invariant μ term is $(H_1^u H_1^d + H_2^u H_2^d)$, and no $H_3^u H_3^d$ and no mixing between the Q_6 doublet and singlet Higgs multiplets are allowed. Therefore, there is an accidental global $SU(2)$, implying the existence of Nambu-Goldstone modes. In [44] the Higgs sector is extended to include a certain set of SM singlet Higgs multiplets to avoid this problem. It has been further found that spontaneous breaking of the flavor symmetry Q_6 as well as CP invariance can be realized without breaking the accidental Z_2 invariance (28). It has been also shown that, although the scale of the singlet sector is of the same order as the heavy $SU(2)$ doublet Higgs bosons, an effective μ term W^{eff} along with the soft-supersymmetry-breaking Lagrangian $\mathcal{L}_{\text{soft}}^{\text{eff}}$ can well describe the original theory in that sector, where

$$W^{\text{eff}} = \mu^{++} (H_+^u H_+^d + H_-^u H_-^d) + \mu^{+3} H_+^u H_3^d + \mu^{3+} H_3^u H_+^d, \quad (29)$$

$$\begin{aligned} \mathcal{L}_{\text{soft}}^{\text{eff}} = & m_{H^u}^2 (|h_+^u|^2 + |h_-^u|^2) + m_{H_3^u}^2 |h_3^u|^2 + m_{H^d}^2 (|h_+^d|^2 + |h_-^d|^2) + m_{H_3^d}^2 |h_3^d|^2 \\ & + [B^{++} (h_+^u h_+^d + h_-^u h_-^d) + B^{+3} h_+^u h_3^d + B^{3+} h_3^u h_+^d + h.c.] . \end{aligned} \quad (30)$$

(The notation H_+^u etc should be self-evident, and the A terms are suppressed.) The parameters μ 's and B 's are complex, which come from the complex VEVs of the singlet Higgs fields of the original theory. So, the effective superpotential (29) and the effective soft-supersymmetry-breaking Lagrangian (30) break Q_6 and CP softly. However, thanks to (28), the VEVs of the form

$$\langle \phi_-^{u,d0} \rangle = 0, \quad \langle \phi_+^{u,d0} \rangle = \frac{v_+^{u,d}}{\sqrt{2}} \exp i\theta_+^{u,d}, \quad \langle \phi_3^{u,d0} \rangle = \frac{v_3^{u,d}}{\sqrt{2}} \exp i\theta_3^{u,d} \quad (31)$$

can be realized. (See (14) for the notation.)

The CKM mixing matrix is given by

$$V_{\text{CKM}} = (U_L^u)^\dagger U_L^d = O_L^{uT} P_q O_L^d, \quad (32)$$

where

$$\begin{aligned} P_q = & \text{diag.} (1, \exp(i2\theta_q), \exp(i\theta_q)), \quad \theta_q = \theta_+^u - \theta_+^d - \theta_3^u + \theta_3^d, \\ O_{uL} \simeq & \begin{pmatrix} 0.9992 & 0.04037 & 9.371 \times 10^{-6} \\ 0.04029 & -0.9974 & 0.05978 \\ -2.422 \times 10^{-3} & 0.05973 & 0.9982 \end{pmatrix}, \end{aligned} \quad (33)$$

$$O_{uR} \simeq \begin{pmatrix} -0.9992 & 0.04037 & -9.025 \times 10^{-5} \\ 0.04029 & 0.9973 & 0.06207 \\ -2.515 \times 10^{-3} & -0.06202 & 0.9981 \end{pmatrix}, \quad (34)$$

$$O_{dL} \simeq \begin{pmatrix} 0.9760 & -0.2176 & -1.945 \times 10^{-3} \\ -0.2174 & -0.9756 & 0.03116 \\ 8.679 \times 10^{-3} & 0.02999 & 0.9995 \end{pmatrix},$$

$$O_{dR} \simeq \begin{pmatrix} -0.9693 & -0.2460 & 1.330 \times 10^{-4} \\ -0.2190 & 0.8628 & 0.4557 \\ 0.1122 & -0.4416 & 0.8901 \end{pmatrix}.$$

The nine independent theory parameters are $Y_a^{u,d} v_3^{u,d}$, $Y_c^{u,d} v_3^{u,d}$, $Y_b^{u,d} v_+^{u,d}$, $Y_{b'}^{u,d} v_+^{u,d}$ and θ_q , which describe the CKM parameters and the quark masses as mentioned. The set of the theory parameters is thus over-constrained. Therefore, there is not much freedom in the parameter space, and so it is sufficient to consider a single point in the space of the theory parameters of this sector. The orthogonal matrices (34) are obtained for the following Yukawa couplings:

$$\begin{aligned} Y_a^u v_3^u &= 1.409 m_t, Y_c^u v_3^u = 2.135 \times 10^{-4} m_t, Y_b^u v_+^u = 0.0847 m_t, Y_{b'}^u v_+^u = 0.0879 m_t, \\ Y_a^d v_3^d &= 1.258 m_b, Y_c^d v_3^d = -6.037 \times 10^{-3} m_b, Y_b^d v_+^d = 0.0495 m_b, Y_{b'}^d v_+^d = 0.6447 m_b, \\ \theta_q &= -0.7125. \end{aligned} \quad (35)$$

With these parameter values we obtain [45]

$$m_u/m_t = 0.609 \times 10^{-5}, m_c/m_t = 3.73 \times 10^{-3}, m_d/m_b = 0.958 \times 10^{-3}, \quad (36)$$

$$m_s/m_b = 1.69 \times 10^{-2}, |V_{\text{CKM}}| = \begin{pmatrix} 0.9740 & 0.2266 & 0.00361 \\ 0.2264 & 0.9731 & 0.0414 \\ 0.00858 & 0.0407 & 0.9991 \end{pmatrix}, \quad (37)$$

$$|V_{td}/V_{ts}| = 0.211, \sin 2\beta(\phi_1) = 0.695, \bar{\rho} = 0.152, \bar{\eta} = 0.343. \quad (38)$$

The mass ratio (37) is defined at M_Z and consistent with the recent up-dates of [47], and the CKM parameters above agree very well with those of Particle Data Group [48] and CKM fitter groups [2, 3]. (See [46] for the prediction of the model in the lepton sector.)

So far we have discussed only the Yukawa sector. To compute the one-loop corrections to δ_{LL} 's, we need to fix the Higgs sector. It is convenient to make a phase rotation of the Higgs superfields so that their VEVs become real:

$$\tilde{H}_\pm^{u,d} = H_\pm^{u,d} e^{-i\phi_\pm^{u,d}}, \tilde{H}_3^{u,d} = H_3^{u,d} e^{-i\phi_3^{u,d}}. \quad (39)$$

Then we define

$$\begin{pmatrix} \Phi_L^{u,d} \\ \Phi_H^{u,d} \\ \Phi_-^{u,d} \end{pmatrix} := \begin{pmatrix} \cos \gamma^{u,d} & \sin \gamma^{u,d} & 0 \\ -\sin \gamma^{u,d} & \cos \gamma^{u,d} & 0 \\ 0 & 0 & 1 \end{pmatrix} \cdot \begin{pmatrix} \tilde{H}_3^{u,d} \\ \tilde{H}_+^{u,d} \\ \tilde{H}_-^{u,d} \end{pmatrix} \quad (40)$$

where

$$\cos \gamma^{u,d} = v_3^{u,d}/v^{u,d}, \quad \sin \gamma^{u,d} = v_+^{u,d}/v^{u,d}, \quad v^{u,d} = \sqrt{(v_3^u)^2 + (v_+^u)^2}. \quad (41)$$

The components of the $SU(2)$ doublet Higgs multiplets are defined

$$\Phi_I^u = \begin{pmatrix} \Phi_I^{u+} \\ \Phi_I^{u0} \end{pmatrix}, \quad \Phi_I^d = \begin{pmatrix} \Phi_I^{d0} \\ \Phi_I^{d-} \end{pmatrix}, \quad I = L, H, -. \quad (42)$$

The light and heavy MSSM-like Higgs scalars are then given by

$$\begin{aligned} (v + h - iX)/\sqrt{2} &= (\phi_L^{d0})^* \cos \beta + (\phi_L^{u0}) \sin \beta, \\ (H + iA)/\sqrt{2} &= -(\phi_L^{d0})^* \sin \beta + (\phi_L^{u0}) \cos \beta, \\ G^+ &= -(\phi_L^{d-})^* \cos \beta + (\phi_L^{u+}) \sin \beta, \quad H^+ = (\phi_L^{d-})^* \sin \beta + \phi_L^u \cos \beta, \end{aligned} \quad (43)$$

where X and G^+ are the Nambu-Goldstone fields, ϕ 's are scalar components of (42), and $v = \sqrt{v_u^2 + v_d^2} (\simeq 246 \text{ GeV})$ and $\tan \beta = v_u/v_d$.

B. Calculation of δ 's

To apply the general formula (16), (19), (22) and (24), we have to compute the mass matrix for the charged Higgs bosons \mathbf{M}_c^2 and fermions \mathbf{M}_h^F . Note that since the Z_2 invariance is unbroken, Z_2 even and odd fields do not mix with each other. We find that the mass matrix of the charged Z_2 even Higgs bosons has the form

$$\mathbf{M}_{c,\text{even}}^2 = \begin{pmatrix} \frac{2B_L}{s_{2\beta}} + c_W^2 M_Z^2 & -\hat{m}_{uLH}^2/c_\beta & -\hat{m}_{dLH}^{2*}/s_\beta \\ -\hat{m}_{uLH}^{2*}/c_\beta & -\hat{m}_{uH}^2 - c_{2\beta} s_W^2 M_Z^2 & B_H^* \\ -\hat{m}_{dLH}^2/s_\beta & B_H & -\hat{m}_{dH}^2 + c_{2\beta} s_W^2 M_Z^2 \end{pmatrix} \quad (44)$$

in the basis of $(H^+, \phi_H^{u+}, (\phi_H^{d-})^*)$, and that of the the charged Z_2 odd Higgs bosons is

$$\mathbf{M}_{c,\text{odd}}^2 = \begin{pmatrix} -m_{H^u}^2 + |\mu^{++}|^2 + c_{2\beta} c_W^2 M_Z^2/2 & B^{+++} \\ B^{++} & -m_{H^d}^2 + |\mu^{++}|^2 - c_{2\beta} c_W^2 M_Z^2/2 \end{pmatrix} \quad (45)$$

in the $(\phi_-^{u+}, (\phi_-^{d-})^*)$ basis, where $c_W = \cos \theta_W = (1 - 0.23)^{1/2}$, $s_{2\beta} = \sin 2\beta$ etc. The mass parameters in (44) and (45) are defined as

$$\begin{aligned} -\hat{m}_{u(d)LH}^2 &= -(c_{\gamma^{u,d}} s_{\gamma^{u,d}} m_{H^{u(d)}}^2 - m_{H_3^{u(d)}}^2) + \mu_H \mu_{LH(HL)}^* + \mu_{HL(LH)} \mu_L^*, \\ -\hat{m}_{u(d)H}^2 &= -(c_{\gamma^{u(d)}}^2 m_{H^{u(d)}}^2 + s_{\gamma^{u(d)}}^2 m_{H_3^{u(d)}}^2) + |\mu_H|^2 + |\mu_{HL(LH)}|^2, \\ B_L &= s_{\gamma^u} s_{\gamma^d} B^{++} e^{i(\theta_+^u + \theta_+^d)} + s_{\gamma^u} c_{\gamma^d} B^{+3} e^{i(\theta_+^u + \theta_3^d)} + c_{\gamma^u} s_{\gamma^d} B^{3+} e^{i(\theta_3^u + \theta_+^d)}, \\ B_H &= c_{\gamma^u} c_{\gamma^d} B^{++} e^{i(\theta_+^u + \theta_+^d)} - c_{\gamma^u} s_{\gamma^d} B^{+3} e^{i(\theta_+^u + \theta_3^d)} - s_{\gamma^u} c_{\gamma^d} B^{3+} e^{i(\theta_3^u + \theta_+^d)}, \end{aligned} \quad (46)$$

$$\begin{aligned}
\mu_{HL} &= c_{\gamma^u} s_{\gamma^d} \mu^{++} e^{i(\theta_+^u + \theta_+^d)} + c_{\gamma^u} c_{\gamma^d} \mu^{+3} e^{i(\theta_+^u + \theta_3^d)} + -s_{\gamma^u} s_{\gamma^d} \mu^{3+} e^{i(\theta_3^u + \theta_+^d)} , \\
\mu_{LH} &= s_{\gamma^u} c_{\gamma^d} \mu^{++} e^{i(\theta_+^u + \theta_+^d)} - s_{\gamma^u} s_{\gamma^d} \mu^{+3} e^{i(\theta_+^u + \theta_3^d)} + c_{\gamma^u} c_{\gamma^d} \mu^{3+} e^{i(\theta_3^u + \theta_+^d)} , \\
\mu_L &= s_{\gamma^u} s_{\gamma^d} \mu^{++} e^{i(\theta_+^u + \theta_+^d)} + s_{\gamma^u} c_{\gamma^d} \mu^{+3} e^{i(\theta_+^u + \theta_3^d)} + c_{\gamma^u} s_{\gamma^d} \mu^{3+} e^{i(\theta_3^u + \theta_+^d)} , \\
\mu_H &= c_{\gamma^u} c_{\gamma^d} \mu^{++} e^{i(\theta_+^u + \theta_+^d)} - c_{\gamma^u} s_{\gamma^d} \mu^{+3} e^{i(\theta_+^u + \theta_3^d)} - s_{\gamma^u} c_{\gamma^d} \mu^{3+} e^{i(\theta_3^u + \theta_+^d)} .
\end{aligned}$$

For the charginos we find

$$\mathbf{M}_{h,\text{even}}^F = \begin{pmatrix} M_2 & \sqrt{2} c_W s_\beta M_Z & 0 \\ \sqrt{2} c_W c_\beta M_Z & \mu_L & \mu_{HL} \\ 0 & \mu_{LH} & \mu_H \end{pmatrix}, \quad M_{h,\text{odd}}^F = \mu^{++}. \quad (47)$$

$\tan \gamma^u$	-0.1188	$\tan \gamma^d$	-0.9480
$\tan \beta$	3.180		
	[TeV]		[TeV ²]
$\mu^{++} e^{i(\theta_+^u + \theta_+^d)}$	(0.900 + i 0.034)	$B^{++} e^{i(\theta_+^u + \theta_+^d)}$	(4.179 ² + i 0.7094 ²)
$\mu^{+3} e^{i(\theta_+^u + \theta_3^d)}$	(0.230 + i 0.120)	$B^{+3} e^{i(\theta_+^u + \theta_3^d)}$	(2.835 ² + i 0.4286 ²)
$\mu^{3+} e^{i(\theta_3^u + \theta_+^d)}$	(-0.660 + i 0.050)	$B^{3+} e^{i(\theta_3^u + \theta_+^d)}$	(-3.106 ² + i 0.1918 ²)
	[TeV ²]		[TeV ²]
$m_{H^u}^2$	-6.457 ²	$m_{H^d}^2$	-7.265 ²
$m_{H_3^u}^2$	-1.261 ²	$m_{H_3^d}^2$	-2.071 ²

TABLE II: A representative set of parameter values. The phases are not fixed, except for θ_q (see (35) and (33).)

To explicitly calculate δ 's, we consider a representative set of parameter values which is given in Table II. For the parameters given in Table II we find the Higgs mass spectrum:

$$\begin{aligned}
\mathbf{M}_{c,\text{even}} &: (8.25, 5.69, 1.61) \text{ TeV}, \quad \mathbf{M}_{c,\text{odd}} : (8.15, 5.45) \text{ TeV}, \\
\mathbf{M}_{h,\text{even}}^F &: (1.14, 0.506, 0.144) \text{ TeV}, \quad M_{h,\text{odd}}^F = 0.900 \text{ TeV}.
\end{aligned} \quad (48)$$

The lightest one of $\mathbf{M}_{c,\text{even}}$ and two of $\mathbf{M}_{h,\text{even}}^F$ correspond to the charged Higgs boson and fermions of the MSSM. Note that the charged fermions are much lighter than the charged bosons, because the EDM constraints require small μ 's as we see from (16) [11]. We have chosen the parameter values in this way, because we have to suppress three-level FCNCs as well as the EDMs.

The A terms (which are suppressed in the soft-supersymmetry-breaking Lagrangian (30)) and soft scalar mass terms have the same family symmetry as the Yukawa sector even in the effective theory [44]. Consequently, the soft scalar mass matrices have the following

form:

$$\begin{aligned}
\tilde{\mathbf{m}}_{aLL}^2 &= m_{\tilde{a}}^2 \text{diag.} (a_L^a, a_L^a, b_L^a) \quad (a = q, l), \\
\tilde{\mathbf{m}}_{aRR}^2 &= m_{\tilde{a}}^2 \text{diag.} (a_R^a, a_R^a, b_R^a) \quad (a = u, d, e), \\
(\tilde{\mathbf{m}}_{aLR}^2)_{ij} &= A_{ij}^a (\mathbf{m}^a)_{ij} = \tilde{A}_{ij}^a m_{\tilde{a}} (\mathbf{m}^a)_{ij} \quad (a = u, d, e),
\end{aligned} \tag{49}$$

where $m_{\tilde{a}}$ denote the average of the squark and slepton masses, respectively, $(a_{L(R)}^a, b_{L(R)}^a)$ are dimensionless free real parameters, A_{ij}^a are free parameters of dimension one, and \mathbf{m}^a are the respective fermion mass matrices. According to [12, 13] we define the tree-level supersymmetry-breaking soft mass insertions as

$$\delta_{LL(RR)}^{a0} = U_{aL(R)}^\dagger \tilde{\mathbf{m}}_{aLL(RR)}^2 U_{aL(R)}/m_{\tilde{a}}^2, \quad \delta_{LR}^{a0} = U_{aL}^\dagger \tilde{\mathbf{m}}_{aLR}^2 U_{aR}/m_{\tilde{a}}^2 \tag{50}$$

in the super CKM basis. Only the A term contributions are included into the left-right mass matrices $\tilde{\mathbf{m}}_{aLR}^2$ in (49), and the μ term contributions (16) will be added to δ_{LR} separately. Note that $a_{L,R}^a$ and A_{ij}^a are all real, because of CP invariance of the original theory, and that the structure (49) is the consequence of the flavor symmetry Q_6 . Since Q_6 is only spontaneously broken in the original theory and only softly broken by the μ terms and B terms in the effective theory described by (29) with (30), there will be no divergent contributions to the non-diagonal elements of δ_{RR} and δ_{LL} . We have explicitly checked the cancellation of the divergences up to terms proportional to the square of the quark masses times gauge coupling squared, which we have anyway neglected. Since these terms are partially included in the following calculations, the cancellation of the renormalization scale Q dependence is not exact. We find about 0.4% change of the non-diagonal elements of δ_{LL} against the change of Q by two orders of magnitude. In the following calculations we set Q equal to $m_{\tilde{d}}$.

We find from (19), (22) and (24)

$$\begin{aligned}
(\delta_{12}^d)_{LL} &= (\delta_{21}^d)_{LL}^* \simeq -2.6 \times 10^{-4} \Delta a_L^q + (1.2 \times 10^{-3} - i 1.2 \times 10^{-6}) \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right]^2, \\
(\delta_{13}^d)_{LL} &= (\delta_{31}^d)_{LL}^* \simeq -8.7 \times 10^{-3} \Delta a_L^q + (0.50 - i 1.9) \times 10^{-2} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right]^2, \\
(\delta_{23}^d)_{LL} &= (\delta_{32}^d)_{LL}^* \simeq -3.0 \times 10^{-2} \Delta a_L^q - (0.28 + i 8.7) \times 10^{-2} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right]^2, \\
(\delta_{12}^d)_{RR} &= (\delta_{21}^d)_{RR}^* \simeq 5.0 \times 10^{-2} \Delta a_R^d, \quad (\delta_{13}^d)_{RR} = (\delta_{31}^d)_{RR}^* \simeq -0.10 \Delta a_R^d, \\
(\delta_{23}^d)_{RR} &= (\delta_{32}^d)_{RR}^* \simeq 0.39 \Delta a_R^d, \\
\Delta a_L^q &= a_L^q - b_L^q, \quad \Delta a_R^d = a_R^d - b_R^d,
\end{aligned} \tag{51}$$

where terms proportional to Δa_L^q and Δa_R^d are the tree-level insertions. We see that the one-loop effects, their real as well as their imaginary parts, to δ_{LL} are comparable to the tree-level ones, while the imaginary part of the (1, 2) element of the one-loop effect is much

smaller than its real part. This is a good news, because CP violation in the first generation of quarks is very small, while CP violation in the third generation may be large. This is a consequence of the hierarchical structure of the Yukawa couplings as one can see from (27) and (35). We have not included the one-loop corrections to δ_{RR} , because due the smallness of the Yukawa couplings in the down-quark sector they are very small compared with the tree-level contributions. For the left-right insertions we find

$$\begin{aligned}
(\delta_{12}^d)_{LR} &\simeq 1.9(\tilde{A}_1^d - \tilde{A}_2^d) \times 10^{-5} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right] + (2.9 - i \ 0.11) \times 10^{-4} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right]^2, \\
(\delta_{21}^d)_{LR} &\simeq (-2.2\tilde{A}_1^d + 1.7\tilde{A}_2^d) \times 10^{-5} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right] - (2.9 - i \ 0.11) \times 10^{-4} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right]^2, \\
(\delta_{13}^d)_{LR} &\simeq (1.0\tilde{A}_1^{'d} + 4.0\tilde{A}_2^{'d}) \times 10^{-5} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right] - (2.1 - i \ 0.09) \times 10^{-4} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right]^2, \\
(\delta_{31}^d)_{LR} &\simeq 5.8\tilde{A}_2^d \times 10^{-4} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right] - (4.2 - i \ 0.17) \times 10^{-3} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right]^2, \\
(\delta_{23}^d)_{LR} &\simeq 1.4\tilde{A}_2^{'d} \times 10^{-4} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right] - (1.0 - i \ 0.04) \times 10^{-3} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right]^2, \\
(\delta_{32}^d)_{LR} &\simeq -2.3\tilde{A}_2^d \times 10^{-2} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right] + (1.7 - i \ 0.07) \times 10^{-2} \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right]^2,
\end{aligned} \tag{53}$$

where the first terms come from the A terms, and the second ones are from the μ terms (16), and \tilde{A}^d 's in (53) are dimensionless free parameters. The imaginary part of the $(1, 1)$ element of δ_{LR} is strongly constraint by EDMs [18, 20]. The A -term contributions are real (because of the CP invariance of the original theory), while the μ term contributions are complex (because spontaneous CP violation generates complex μ terms). We find, using (16),

$$\begin{cases} \text{Im}(\delta_{11}^d)_{LR} = 3.7 \\ \text{Im}(\delta_{11}^u)_{LR} = 1.3 \end{cases} \times 10^{-6} \times \left[\frac{0.5 \text{ TeV}}{m_{\tilde{d}}} \right]^2, \tag{54}$$

which are of the order of the upper bound [18].

C. Mixing of the neutral mesons and new CP phases

We now apply our results (51) and (53) to the mixing of the neutral meson systems. Here we assume that the tree-level contributions to the mixing coming from the heavy neutral Higgs boson exchange are small. In [10, 11, 27] it has been found that if their masses are larger than several TeV in the present model, then FCNCs and CP are suppressed. In the present case with the parameters given in Table II, the mass of the lightest flavor-changing neutral Higgs boson is 5.7 TeV, and so the assumption may be justified. The total matrix

element M_{12}^q in the neutral meson mixing can be written as

$$M_{12}^q = M_{12}^{SM,q} + M_{12}^{SUSY,q} , \quad (55)$$

where $M_{12}^{SM,q}$ is the SM contribution, and $M_{12}^{SUSY,q}$ is the SUSY contribution, whose dominant contribution is given by [18] (see e.g. [49] for a more refined calculation)

$$\begin{aligned} M_{12}^{SUSY,s} = & -\frac{\alpha_S^2}{324m_{\tilde{d}}^2} M_s f_{B_s}^2 B_s \left\{ [(\delta_{32}^d)_{LL}^2 + (\delta_{32}^d)_{RR}^2] [24x f_6(x) + 66\tilde{f}_6(x)] \right. \\ & + (\delta_{32}^d)_{LL} (\delta_{32}^d)_{RR} [(384 R_s + 72) x f_6(x) - (24R_s - 36) \tilde{f}_6(x)] \\ & - 132 [(\delta_{32}^d)_{LR}^2 + (\delta_{23}^d)_{LR}^{*2}] R_s x f_6(x) \\ & \left. - (\delta_{32}^d)_{LR} (\delta_{23}^d)_{LR}^* [144 R_s + 84] \tilde{f}_6(x) \right\} , \end{aligned} \quad (56)$$

where

$$\begin{aligned} R_s &= \left(\frac{M_s}{m_s + m_d} \right)^2 , \\ f_6(x) &= \frac{6(1+3x) \ln x + x^3 - 9x^2 - 9x + 17}{6(x-1)^5} , \\ \tilde{f}_6(x) &= \frac{6x(1+x) \ln x - x^3 - 9x^2 + 9x + 1}{3(x-1)^5} \quad \text{with } x = m_{\tilde{g}}/m_{\tilde{d}} , \end{aligned} \quad (57)$$

similarly for K and B_d , and $m_{\tilde{g}}$ is the gluino mass. For the calculations below we assume that the bag parameters B_K, B_d, B_s are one, and $\alpha_S = 0.12$. The other parameters are given in Table III. Since (56) is a one-loop result, the one-loop effect to the soft mass insertions in fact means a two-loop effect like Fig. 3

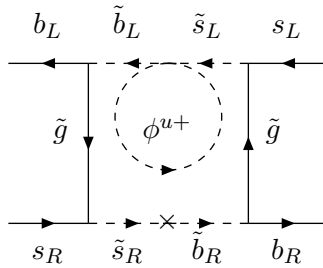


FIG. 3: An example of two-loop contribution to $M_{12}^{SUSY,s}$. One-loop contribution to the insertion $(\delta_{32}^d)_{LL}$, the ϕ^{u+} loop in the box, means a two-loop effect on $M_{12}^{SUSY,s}$.

We follow [7] to parameterize new physics effects as (see (3))

$$M_{12}^{SM,q} + M_{12}^{SUSY,q} = M_{12}^{SM,q} \cdot \Delta_q , \quad (58)$$

Input		Input	
f_K	$(159.8 \pm 1.4 \pm 0.44) \times 10^{-3} \text{ GeV}$	f_{B_d}	$0.194 \pm 0.032 \text{ GeV}$
f_{B_s}	$0.240 \pm 0.040 \text{ GeV}$		
M_K	$0.497648 \pm 0.000022 \text{ GeV}$	ΔM_K^{exp}	$(0.5292 \pm 0.0009) \times 10^{-2} \text{ ps}^{-1}$
M_s	$5.3661 \pm 0.0006 \text{ GeV}$	ΔM_s^{exp}	$17.77 \pm 0.10 \pm 0.07 \text{ ps}^{-1}$
M_d	$5.27950 \pm 0.00033 \text{ GeV}$	ΔM_d^{exp}	$0.507 \pm 0.005 \text{ ps}^{-1}$
$m_d(2\text{GeV})$	$(5.04^{+0.96}_{-1.54}) \times 10^{-3} \text{ GeV}$	$m_s(2\text{GeV})$	$0.105^{+0.025}_{-0.035} \text{ GeV}$
$m_d(m_b)$	$(4.23^{+1.74}_{-1.71}) \times 10^{-3} \text{ GeV}$	$m_s(m_b)$	$0.080 \pm 0.022 \text{ GeV}$
$m_b(m_b)$	$4.20 \pm 0.07 \text{ GeV}$		

TABLE III: Parameter values used in the text (see also Ref. [50]). For the calculations in the text we use only the central values. $f_K, M_{K,d,s}, \Delta M_{K,d,s}^{\text{exp}}$ are from [48]. f_{B_s} belongs to the conservative sets of [7] (see the references therein), and f_{B_d} is obtained from f_{B_s}/ξ with $\xi = 1.24$. $m_d(2\text{GeV})$ and $m_s(2\text{GeV})$ are from [48], while those at m_b are taken from [47].

and consider the observables ΔM_q , $\Delta \Gamma_q$ and the flavor specific CP-asymmetry a_{sl}^q in terms of the complex number $\Delta_q = |\Delta_q| e^{i\phi_q^\Delta}$, where $q = d, s$, and

$$\begin{aligned} \Delta M_q &= 2|M_{12}^{SM,q}| \cdot |\Delta_q|, \quad \Delta \Gamma_q = 2|\Gamma_{12}^q| \cos(\phi_q^{SM} + \phi_q^\Delta), \\ a_{sl}^q &= \frac{|\Gamma_{12}^q|}{|M_{12}^{SM,q}|} \cdot \frac{\sin(\phi_q^{SM} + \phi_q^\Delta)}{|\Delta_q|}. \end{aligned} \quad (59)$$

The SM values are given e.g. in [7], in which the results of [30–34] are used:

$$\begin{aligned} 2 M_{12}^{SM,d} &= 0.56(1 \pm 0.45) \exp(i0.77) \text{ ps}^{-1}, \\ 2 M_{12}^{SM,s} &= 20.1(1 \pm 0.40) \exp(-i0.035) \text{ ps}^{-1}, \\ \Delta \Gamma_d^{SM} &= (26.7^{+5.8}_{-6.5}) \times 10^{-4} \text{ ps}^{-1}, \quad \Delta \Gamma_s^{SM} = 0.096 \pm 0.039 \text{ ps}^{-1}, \end{aligned} \quad (60)$$

where the errors are dominated by the uncertainty in the decay constants and bag parameters.

We use the central values of (7), (8) and Table III for our calculations, while requiring the constraints

$$\left. \begin{matrix} 0.6 \\ 0.8 \end{matrix} \right\} < \frac{\Delta M_{d,s}}{\Delta M_{d,s}^{\text{exp}}} < \left\{ \begin{matrix} 1.4 \\ 1.2 \end{matrix} \right., \quad \frac{2|M_{12}^{SUSY,K}|}{\Delta M_K^{\text{exp}}} < \left\{ \begin{matrix} 2 & I \\ 1 & II \end{matrix} \right. \quad (61)$$

and

$$\frac{\text{Im} M_{12}^{SUSY,K}}{\sqrt{2} \Delta M_K^{\text{exp}}} < \epsilon_K = 2.2 \times 10^{-3}, \quad (62)$$

where I and II are a conservative and an optimistic set of constraints, respectively.

As we can see from (53), except for $(\delta_{23,32}^d)_{LR}$, the μ term contributions (the second terms) are larger than the A term contributions by an order of magnitude (if A 's are of $O(1)$). Therefore, we include only the μ term contributions to $M_{12}^{SUSY,K}$ and $M_{12}^{SUSY,d}$. Further, $(\delta_{23,32}^d)_{LR}$ are constrained by $b \rightarrow s\gamma$ [16, 18, 19], and have to satisfy $|(\delta_{23,32}^d)_{LR}| \lesssim 10^{-2}$. So, $|(\delta_{23,32}^d)_{LR}|$ are saturated with the μ term contribution. We could choose a positive $O(1)$ value for \tilde{A}_2^d so that the A term contribution cancels the μ term contribution, but the effect on $|(\delta_{23,32}^d)_{LR}|$ is negligibly small. We therefore neglect the A term contribution in $(\delta_{23,32}^d)_{LR}$, too. Under this situation, as we can see from (51), (53) and (56), given the set of parameters of Table II and III with $x = m_{\tilde{g}}/m_{\tilde{d}} = 1$, the free parameters are only Δa_L^d and Δa_R^d . So, it is absolutely non-trivial to satisfy the constraints (61) and (62) while having a large CP violation in the B^0 mixing.

In Fig. 4 we plot ϕ_d^Δ against ϕ_s^Δ for the parameter values given in Tables II and III. The green (red) region satisfies the constraints I (II) with (62). The SM value, (7) and (8), is denoted by \bullet . Since no large CP phase can be generated without the loop effects of the extra heavy Higgs bosons in this model [10, 11], the large $\phi_{d,s}^\Delta$ in the predicted region of Fig. 4 is entirely due the loop effects. Fig. 5 shows the theoretical values in the $\phi_s^\Delta - A_{sl}^b$ plane, where the same sign dimuon asymmetry A_{sl}^b is computed from (9). Here we have imposed only the conservative constraint I of (61) with (62), and assumed that $|\Gamma_{12}^d|/|M_{12}^{SM,d}| = (52.6^{+11.5}_{-12.8}) \times 10^{-4}$, $|\Gamma_{12}^s|/|M_{12}^{SM,s}| = (4.97 \pm 0.94) \times 10^{-3}$ [7]. The SM value \bullet is obtained from (2) and (8). The theoretical values in Fig. 5 should be compared with the recent experimental measurements; $A_{sl}^b = -(9.57 \pm 2.51 \pm 1.46) \cdot 10^{-3}$ [6] and $\phi_s^\Delta \in [-0.2, -2.8]$ [41] ($\phi_s^\Delta \in [0, -1]$ [42] at 68% C.F.). As we see from Fig. 5, the theoretical value for ϕ_s^Δ is consistent with the observation, but that of A_{sl}^b is still at least a factor of 5 smaller than the D0 measurement of the dimuon asymmetry [6].

V. CONCLUSION

In this paper we have considered the Higgs sector of a generic supersymmetric extension of the SM, while assuming that there are more than one pair of $SU(2)_L$ doublet Higgs supermultiplets. Such a case is realized in models with a low-energy flavor symmetry. We have calculated the one-loop effects of the extra Higgs multiplets to the soft mass insertions. We have found that under a certain circumstance the loop effects can give rise to large contributions to the soft mass insertions, which means that they can generate large FCNCs and CP violations. We have applied a supersymmetric extension of the SM based on the discrete Q_6 family symmetry. Due to the flavor symmetry, the flavor-non-diagonal loop contributions are finite in this model. We have calculated the supersymmetric contribution to the non-diagonal matrix element M_{12} of the neutral meson systems. In particular, we have calculated the extra phases $\phi_{d,s}^\Delta$ and the flavor-specific CP-asymmetries $a_{sl}^{d,s}$ in the B^0

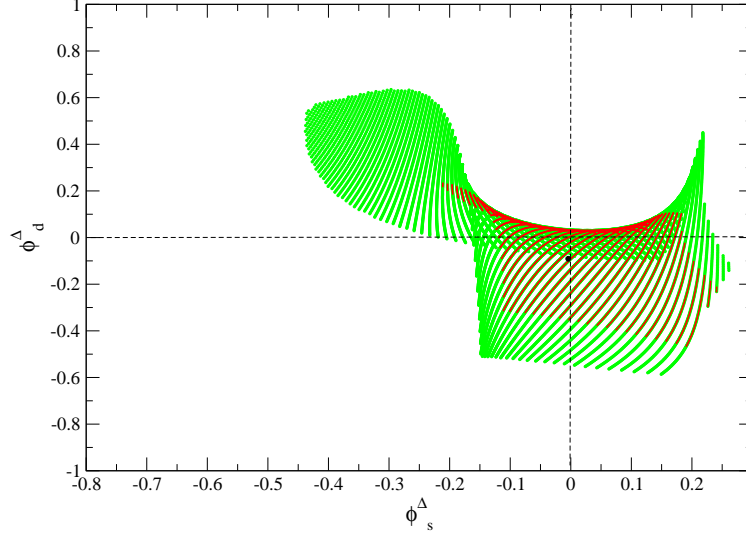


FIG. 4: The theoretical values in the $\phi_s^\Delta - \phi_d^\Delta$ plane for the parameter values given in Tables II and III. The constraint I (II) of (61) with (62) is satisfied by the green (red) region. The \bullet is the SM value.

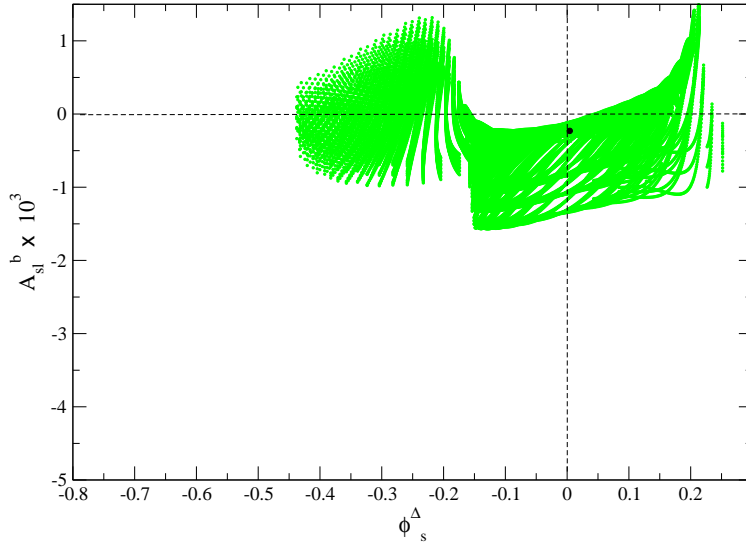


FIG. 5: The theoretical values in the $\phi_s^\Delta - A_{sl}^b$ plane for the same parameter values as Fig. 4. We impose only the constraint I with (62). The SM value \bullet is obtained from (2) and (8). The recent experimental values are, respectively, $A_{sl}^b = -(9.57 \pm 2.51 \pm 1.46) \cdot 10^{-3}$ [6] and $\phi_s^\Delta \in [-0.2, -2.8]$ [41] ($\phi_s^\Delta \in [0, -1]$ [42]).

mixing and that the value of β_s of the model is consistent with the recent CDF measurement [42]. As for the same sign dimuon asymmetry A_{sl}^b we obtain values which are one order of magnitude larger than the SM model value. Nevertheless, they are at least a factor of 5 smaller than the D0 measurement of the dimuon asymmetry [6].

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